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Chiral Fermion Delocalization in Deconstructed Higgsless Theories

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Abstract

I construct a renormalizable $SU(2)^{89} \times U(1)$ gauge theory with standard-model-like phenomenology for the gauge bosons masses and the weak interactions of the light fermions (including the b) but in which all vacuum expectation values are about 2 TeV. This is a deconstructed version of a Higgsless model with a flat extra dimension. The fermions are delocalized on the theory space in an unusual way, with LH and RD fermions on alternate nodes.

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In this note, I report on an exercise in gauge theory model-building that could have been done thirty years ago. The model is a conventional renormalizable[1] quantum field theory with gauge symmetry spontaneously broken by the vacuum expectation values of elementary scalar fields. The model is designed to reproduce the electroweak interactions of the standard model to a good approximation in tree approximation. The obvious difference between this model and the conventional standard model with an elementary Higgs boson is the size of the gauge group and the fermion representation. The electroweak gauge group is $SU(2)^{89} \times U(1)$ and the number of fermion representations is similarly swollen. The reason I consider this ludicrously large structure is to illustrate in a very explicit way how a so-called Higgsless model [1, 2] can very nearly reproduce the low-energy phenomenology of the standard model. In particular, the particle masses and the couplings of the light fermions are within a percent or two of their standard model values. Of course this model is not without elementary Higgs bosons. However, it is a deconstructed [3, 4] version of a Higgsless model [5, 6, 7, 8, 9, 10, 11, 12, 13, 14], and the translation of "Higgslessness" into the language of 70s model building is simply that the vacuum expectation values that break the $SU(2)$ symmetries are all much larger than the 250 GeV of the standard model. In the explicit example I will describe, the VEVs are all greater than 2 TeV. I find this rather remarkable and I simply could not have imagined back in the 70s that such a thing was possible.

It is quite clear in retrospect why models like this were not constructed in the 70's. From the 4-dimensional point of view, these models look crazy, with hundreds of apparently extra structures and parameters! What one gains from thinking about extra dimensions is motivation to consider models that look very complicated from the 4-d point of view but which, from the 5-d point of view, are really rather simple. Indeed, in the explicit model I discuss here, I will make an assumption related to translation invariance in the extra dimension that reduces the number of parameters to just three more than in the standard model. For me, this is all that extra dimensions are good for. I try to deconstruct them as quickly as possible so that I know what I am doing. But I think perhaps that the youngsters who have been weaned on a geometrical picture gain some intuition directly from 5-dimensional thinking.

What is new in this note, I believe, is a more explicit and different treatment of what is called in the 5-d language the "delocalization" of the fermions. It is necessary to spread the the fermions out in the extra dimensions to produce a negative contribution to the S parameter [15, 16, 17, 14, 18]. In the 4-d language of this paper, this simply means that the light fermion doublets are linear combinations of fields that transform under different $SU(2)$ subalgebras. I will introduce a very efficient chiral fermion delocalization in which the left-handed and right-handed fermions appear on alternate nodes. I find this a bit confusing from the 5-d point of view. More of this later.

I will begin by describing the model without much further motivation. Once we have all the pieces and I have described a particular choice of parameters that does what I want, I will step back and discuss things more generally.

[1] But see the discussion in footnote 2 on page 3.

The gauge structure of the model is summarized in pictorial form in figure 1. There are

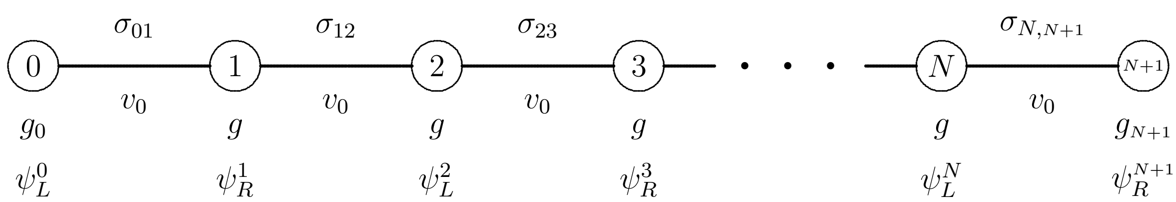

Figure 1: The Moose diagram associated with the example.

$N+1$ $SU(2)$ factors in the gauge group, and N is constrained to be even, but is otherwise free. The links and nodes of the figure form a "theory space" that corresponds via deconstruction to the configuration space of the extra dimension. Nodes 1 through N are associated with the "bulk" of the extra dimension, between the ends of the figure which correspond to two 4-d "branes." I have assumed that all the bulk groups have the same gauge coupling, consistent with translation invariance in a flat extra dimension.[2] The links in the figure are associated with 2×2 real σ fields satisfying

$$\sigma = s + i\,\vec{\tau} \cdot \vec{p} \qquad \sigma^*_{j,j+1} = \tau_2 \sigma_{j,j+1} \tau_2 \tag{1}$$

each with vacuum expectation value v_0, again preserving the translation invariance

$$\langle \sigma_{j,j+1} \rangle = v_0 \tag{2}$$

Note that the VEVs can all be rotated to be proportional to the identity.

Because of the assumption of translation invariance, the gauge sector of the model is determined by only four parameters, v_0, g_0, g, and g_{N+1}, compared to three in the standard model, v, e and $\sin\theta$. Though it will not play any important role in the tree-level analysis, it is reasonable to assume a translation invariant set of potentials for the scalars as well.

To have fermion delocalization, we need to spread the fermion doublets from one end of the theory space in figure 1 to the other. The most efficient way to do this, I believe, is to alternate with the LH doublets on the even nodes and RH doublets on the odd nodes, as shown in figure 1. This setup has a number of nice features, which we will discuss below. This only works for even N, which is why we have assumed that N is even.

The field ψ_R^{N+1} is written as a doublet, but because the $N+1$ group is a $U(1)$, the top and bottom components are independent.

$$\psi_R^{N+1} = \begin{pmatrix} U_R \\ D_R \end{pmatrix} \tag{3}$$

[2] Technically, this is not a natural constraint. The breakdown of translation invariance at the edges will, in high enough order, produce infinite j-dependent renormalizations of the bulk couplings, but nobody has worried about that for years.

If we write the fields in a column vector in the theory space where the LH and RH fields alternate on the even and odd components,[3]

$$[\Psi]_{2j} = \psi_L^{2j} \,, \quad [\Psi]_{2j+1} = \psi_R^{2j+1} \;\text{ for } j = 0 \text{ to } N/2. \tag{4}$$

and the Yukawa couplings can be written as

$$\begin{gathered} \overline{\Psi}\,\mathcal{A}^f\,\Psi \quad \text{with} \quad [\mathcal{A}^f]_{j,k} = 0 \;\text{ if } j \neq k \pm 1 \\ [\mathcal{A}^f]_{j,j+1} = a_{j,j+1}\sigma_{j,j+1} \,, \;\; [\mathcal{A}^f]_{j+1,j} = a_{j,j+1}\sigma^\dagger_{j,j+1} \;\text{ for } j = 0 \text{ to } N \end{gathered} \tag{5}$$

Because we have imposed translation invariance in the bulk for the gauge couplings and VEVs, we should also impose translation invariance for the Yukawa coupling. This, however, is somewhat peculiar given our fermion representation, because of the alternation of LH and RH fermions in (4). Thus we want our discrete translation invariance in the bulk to have the form

$$j \to j+1 \qquad L \leftrightarrow R \tag{6}$$

changing parity!. This seems a bit odd, but it has an important consequence. Because of (6), the fermion mass matrix in the bulk has the simple form

$$\overline{\Psi}\,\mathcal{M}\,\Psi \;\text{ with }\; [\mathcal{M}]_{j,k} = 0 \;\text{ if } j \neq k \pm 1 \qquad [\mathcal{M}]_{j,j+1} = a v_0 \,, \;\; [\mathcal{M}]_{j+1,j} = a v_0 \tag{7}$$

which automatically has zero modes that are spread over the whole bulk. I will say more about this below.

We will also assume that the form (7) is flavor independent, and in fact that the only flavor dependence is on the $U(1)$ brane at $j = N+1$. We will also break the translation invariance on the brane at $j = 0$. We will need this freedom below. But we will assume that the coupling at the $j = 0$ brane is flavor independent. It might be possible to get away with very small differences in the “bulk” couplings of the various flavors and the coupling to the $j = 0$ brane, but this is a very dangerous path, likely to exacerbate the problem of universality violation and flavor changing neutral currents, and I want to see how far we can get without walking this plank.

Making these assumptions, the Yukawa couplings become

$$\begin{gathered} \overline{\Psi}\,\mathcal{A}^f\,\Psi \quad \text{with} \quad [\mathcal{A}^f]_{j,k} = 0 \;\text{ if } j \neq k \pm 1 \\ [\mathcal{A}^f]_{j,j+1} = a\,\sigma_{j,j+1} \,, \;\; [\mathcal{A}^f]_{j+1,j} = a\,\sigma^\dagger_{j,j+1} \;\text{ for } j = 1 \text{ to } N-1 \\ [\mathcal{A}^f]_{01} = \epsilon_0 a\,\sigma_{01} \,, \;\; [\mathcal{A}^f]_{10} = \epsilon_0 a\,\sigma^\dagger_{01} \\ [\mathcal{A}^f]_{N,N+1} = \epsilon_N^{f_U/f_D} a\,\sigma_{N,N+1} \,, \;\; [\mathcal{A}^f]_{N+1,N} = \epsilon_N^{f_U/f_D} a\,\sigma^\dagger_{N,N+1} \end{gathered} \tag{8}$$

[3]For thirty years I have taught my students not to mix LH and RH fields. Here, for once, it is appropriate, because the form of the fermion mass matrix ensures that the couplings have the right form. There is still something odd about the notation though. Because the fermion states break up into LH states on the even nodes and RH states on the odd, the eigenvalues come in degenerate pairs and there is a superselection rule — there is no need ever to superpose even-node states with odd-node states.

with all flavor dependence in the constant $\epsilon_N^{f_U/f_D}$.

The mass matrix then has the form

$$\begin{array}{c} \overline{\Psi}\,\mathcal{M}\,\Psi \text{ with } [\mathcal{M}]_{j,k}=0 \text{ if } j\neq k\pm 1 \qquad [\mathcal{M}]_{j,j+1}=a\,v_0\,,\ \ [\mathcal{M}]_{j+1,j}=a\,v_0 \\ [\mathcal{M}]_{01}=\epsilon_0 a\,v_0\,,\ \ [\mathcal{M}]_{10}=\epsilon_0 a\,v_0\,,\ \ [\mathcal{M}]_{N,N+1}=\epsilon_N^{f_U/f_D}a\,v_0\,,\ \ [\mathcal{M}]_{N+1,N}=\epsilon_N^{f_U/f_D}a\,v_0 \end{array} \tag{9}$$

As expected, the eigenvectors of this mass matrix come in pairs with degenerate eigenvalues, eigenvectors e_L^j with only even components and e_R^j with only odd components.

Note that the only place where there is a difference between the couplings of the top and bottom of the RH doublet ψ_R^{N+1} is in the term proportional to $\epsilon_N^{f_U/f_D}$. This is always true whatever other assumptions we make about the form of the Yukawa couplings.

For the light quarks and all the leptons, the $\epsilon_N^{f_U/f_D}$ constants will be very small, and to first approximation, we can simply assume

$$\epsilon_N^{\text{light fermions}}\approx 0 \tag{10}$$

In this limit, the eigenvectors for the LH light modes are very simple. For approximately zero eigenvalue, and they all have the form e_L^0 where

$$[e_L^0]_0=\sqrt{\frac{1}{1+N\epsilon_0^2/2}} \qquad [e_L^0]_{2j}=\alpha_j=(-1)^j\,\epsilon_0\,[e_L^0]_0 \qquad [e_L^0]_{2j-1}=0 \tag{11}$$

for $j=1$ to $N/2$. The reader should now begin to see why we need the extra freedom of assuming that

$$\epsilon_0\neq 1 \tag{12}$$

The quantity ϵ_0 determines the amount of delocalization of the LH fermions. In the limit we are considering, we want to keep the delocalization relatively small so we are interested in the region

$$|\epsilon_0|\ll 1 \tag{13}$$

In this limit, the RH light mode is not delocalized. It is stuck on the "brane" – the $N+1$ node

$$[e_R^0]_j=\delta_{j,N+1} \tag{14}$$

For the t quark, however, ϵ_N is not small, the RH mode is delocalized and the LH mode is more complicated than (11). The LH eigenvector e_L^t for the lightest t mode has the form

$$\begin{array}{c} [e_L^t]_{2j-1}=0 \qquad [e_L^t]_{2j}=(-1)^j\,[e_L^t]_0\times \\ \dfrac{(2-\epsilon_0^2-2(1-\epsilon_0^2)\cos c_t)\cos(jc_t+c_t)-(\epsilon_0^2-2(1-\epsilon_0^2)\cos c_t)\sin(jc_t+c_t)\tan(c_t/2)}{\epsilon_0 a} \end{array} \tag{15}$$

for $j=1$ to $N/2$, where

$$c_t=2\arcsin\frac{m_t}{2a\,v_0} \tag{16}$$

This satisfies the eigenvalue equation if $(\epsilon_N^t)^2$ is given by

$$\frac{\sin\frac{c_t}{2}\left(8\cos\frac{(N+1)c_t}{2}\sin\frac{c_t}{2}+4\epsilon_0^2\sin\frac{Nc_t}{2}\right)}{2\epsilon_0^2\cos\frac{(N-1)c_t}{2}-4\sin\frac{c_t}{2}\sin\frac{Nc_t}{2}} \tag{17}$$

Thus we can choose any ϵ_0 and N so long as (17) gives a positive value for $(\epsilon_N^t)^2$.

The mass matrix for the charged gauge bosons is

$$M_c^2 = \frac{1}{4}\widetilde{G}\,\widetilde{V}\,\widetilde{G} \tag{18}$$

where $\widetilde{G}$ is the $(N+1)\times(N+1)$ diagonal matrix of gauge couplings without g_{N+1}

$$[\widetilde{G}]_{jk} = 0 \;\text{ for } j \neq k\,, \quad [\widetilde{G}]_{00} = g_0\,, \quad [\widetilde{G}]_{jj} = g \;\text{ for } j = 1 \text{ to } N\,. \tag{19}$$

and the matrix $\widetilde{V}$ is given by

$$\begin{gathered}[\widetilde{V}]_{jk} = 0 \;\text{ for } j \neq k,\; k\pm 1\,, \quad [\widetilde{V}]_{00} = v_0^2 \quad , \quad [\widetilde{V}]_{jj} = 2v_0^2 \;\text{ for } j = 1 \text{ to } N\,, \\ [\widetilde{V}]_{j,j+1} = [\widetilde{V}]_{j+1,j} = -v_0^2 \;\text{ for } j = 0 \text{ to } N-1\,.\end{gathered} \tag{20}$$

The low energy charged-current weak interactions of the light fermion modes are determined by the inverse of $\widetilde{V}$ and the eigenvector e_L^0:

$$\sqrt{2}\,G_F = \frac{1}{v^2} = \sum_{j,k=0}^{N+1}[\widetilde{V}^{-1}]_{jk}|[e_L^0]_j|^2|[e_L^0]_k|^2 \tag{21}$$

The matrix $\widetilde{V}^{-1}$ satisfies

$$[\widetilde{V}^{-1}]_{jk} = \min(N+1-j, N+1-k)/v_0^2 \tag{22}$$

so

$$\frac{v_0^2}{v^2} = \frac{N+1+\epsilon_0^2N^2/2+\epsilon_0^4N(N^2+2)/12}{(1+\epsilon_0^2N^2/2)^2} \tag{23}$$

The low energy neutral-current weak interactions are then given in terms of $\widetilde{V}^{-1}$ by [19]

$$\sum_{j,k=0}^{N+1}[\widetilde{V}^{-1}]_{jk}\left[T_3|[e_L^0]_j|^2 - \frac{e^2}{g_j^2}Q\right]\,\left[T_3|[e_L^0]_j|^2 - \frac{e^2}{g_k^2}Q\right] \tag{24}$$

The normalization of the T_3^2 term satisfies custodial $SU(2)$ symmetry, so the correction to the ρ parameter is small. The analog of $\sin^2\theta$ as determined by the low energy weak interactions is determined by the coefficient of $T_3\,Q$ in (24) to satisfy

$$\sqrt{2}G_F\sin^2\theta = \sum_{j,k=0}^{N}[\widetilde{V}^{-1}]_{jk}|[e_L^0]_j|^2\,\frac{e^2}{g_k^2} \tag{25}$$

which using (11) and (22), we can write as

$$\frac{(N+1)e^2/g_0^2 + \epsilon_0^2 N^2 e^2/4g_0^2 + N(N+1)e^2/2g^2 + \epsilon_0^2 N(4N^2+3N+2)e^2/24g^2}{1+\epsilon_0^2 N/2} \tag{26}$$

The mass matrix for the neutral gauge bosons is

$$M_n^2 = \frac{1}{4} G\, V\, G \tag{27}$$

where G is the $(N+2)\times(N+2)$ diagonal matrix of gauge couplings

$$[G]_{jk} = 0 \text{ for } j \neq k\,, \quad [G]_{00} = g_0\,, \quad [G]_{jj} = g \text{ for } j = 1 \text{ to } N\,, \quad [G]_{N+1,N+1} = g_{N+1}\,. \tag{28}$$

and the matrix V is

$$\begin{gathered}[V]_{jk} = 0 \text{ for } j \neq k,\, k \pm 1\,, \quad [V]_{jj} = 2v_0^2 \text{ for } j = 1 \text{ to } N\,, \\ [V]_{00} = [V]_{N+1,N+1} = v_0^2 \quad , \quad [V]_{j,j+1} = [V]_{j+1,j} = -v_0^2 \text{ for } j = 0 \text{ to } N\,.\end{gathered} \tag{29}$$

The neutral mass squared matrix given by (27) has, of course, a zero eigenvalue associated with the photon. The photon eigenstate κ_{N+1} is given by

$$[\kappa^{N+1}]_0 = \frac{e}{g_0}\,, \qquad [\kappa^{N+1}]_j = \frac{e}{g} \quad \text{for } j = 1 \text{ to } N, \qquad [\kappa^{N+1}]_{N+1} = \frac{e}{g_{N+1}} \tag{30}$$

Having specified the model in general, let us now look at what happens for a particular set of parameters chosen to produce something like the standard model at low energies. To produce these values, I fixed e, v, M_W, M_Z and m_t, and then scanned over various values of ϵ_0, a and N. To do this efficiently, it is convenient to have analytic expressions for g_0, g_{N+1} and ϵ^t_{N+1} in terms of the fixed values. This is possible because of the translation invariance in bulk, and this and other details of the search procedure will be discussed fully in a forthcoming paper [20]. But it is straightforward, given a set of parameters, N, g_0, g, g_{N+1}, v_0, a, ϵ_0 and ϵ^t_{N+1} to diagonalize the mass matrices numerically and see that it all works.[4] Here is one set of parameters:

$$\begin{gathered}N = 88\,, \quad g = 7.8\,, \quad g_0 = 0.899\,, \quad g_{N+1} = 0.363\,, \\ v_0 = 2.005\,\text{TeV}\,, \quad a = 7\,, \quad \epsilon_0 = 0.1\,, \quad \epsilon^t_{N+1} = 0.309\,.\end{gathered} \tag{31}$$

These give

$$\begin{gathered}\alpha = 1/129\,, \quad v = 250\,\text{GeV}\,, \quad \sin^2\theta = 0.228 \\ M_W = 80.425\,\text{GeV}\,, \quad M_Z = 91.1876\,\text{GeV}\,, \quad m_t = 175\,\text{GeV}\,, \\ \frac{WLL}{\text{s.m.}} = 0.987\,, \quad \frac{Wt_Lb_L}{\text{s.m.}} = 1.052\,, \quad \frac{ZLL}{\text{s.m.}} = 0.985\,, \quad \frac{ZRR}{\text{s.m.}} = 0.972\,, \\ \frac{WWZ}{\text{s.m.}} = 1.097\,, \quad \frac{Zt_Lt_L}{\text{s.m.}} = 1.137\,, \quad \frac{Zt_Rt_R}{\text{s.m.}} = 0.096\,.\end{gathered} \tag{32}$$

[4] Of course, it was less straightforward in the early 70s, when it was a major production to do numerical work on anything more powerful than a slide-rule.

In (32), $\sin^2\theta$ is the tree-level value in the low-energy neutral-current weak interactions, from (25). The couplings of the W and Z are tabulated using the numerical eigenvectors and compared to the tree-level values in the standard model.[5]

In table 1, I tabulate the masses and some of the couplings of the first nine recurrences of the W and Z. Table 2, contains similar information for the first nine recurrences of the light fermions and the t. There is interesting phenomenology here beyond what already

M_W (GeV)	$\frac{WLL}{\text{s.m.}}$	$\frac{Wt_Lb_L}{\text{s.m.}}$	$\frac{W_kWZ}{\text{s.m.}}$	M_Z (GeV)	$\frac{ZLL}{\text{s.m.}}$	$\frac{ZRR}{\text{s.m.}}$	$\frac{Zt_Lt_L}{\text{s.m.}}$	$\frac{Zt_Rt_R}{\text{s.m.}}$	$\frac{Z_kWW}{\text{s.m.}}$
80.42	0.99	1.036	1.097	91.19	0.988	0.973	1.137	0.961	1.097
305.03	0.088	0.239	0.224	309.46	0.099	0.357	0.382	2.871	0.182
567.93	0.224	0.173	0.012	570.47	0.254	0.2	0.197	0.334	0.011
838.58	0.038	0.093	0.008	840.33	0.043	0.137	0.132	1.05	0.007
1111.41	0.116	0.09	0.001	1112.74	0.133	0.104	0.103	0.168	0.001
1384.94	0.023	0.057	0.002	1386.	0.026	0.083	0.079	0.635	0.002
1658.57	0.078	0.06	≈ 0	1659.45	0.089	0.069	0.069	0.109	≈ 0
1931.99	0.017	0.041	0.001	1932.75	0.019	0.059	0.057	0.453	0.001
2205.	0.059	0.045	≈ 0	2205.67	0.068	0.052	0.052	0.078	≈ 0
2477.46	0.013	0.032	≈ 0	2478.05	0.015	0.046	0.044	0.35	≈ 0

Table 1: The W and Z and their first nine recurrences, along with their couplings to quarks compared to the standard model couplings.

appears in the literature [21, 22, 23]. For example, in this class of models, splittings between fermion recurrences are much larger than between gauge boson recurrences because $a > g/2$ and because the chiral delocalization effectively reduces the size of the extra dimension for the fermions in half.[6]

Now for some comments

1. Because we have chosen parameters to give the right masses and couplings for W and Z, we automatically ensure that the S parameter is small and that other low energy tests of the standard model are satisfied. The couplings of the t and the WWZ couplings show more deviation from the tree-level standard model, as one would expect because the effective theory at large energies is very different. Clearly, the values in (31) are finely tuned to produce something that looks like the standard model. On the other hand, this model does not have a lot of parameters to tune. It is surprising, at least to me, that one can do this at all. In particular, we are able to get a large enough t mass without making the Zbb couplings significantly different from those of the other light quarks. This has been a worry in previous works [24, 25, 26].

[5] L (R) stands for any LH (RH) light field.

[6] I am grateful the Sekhar Chivukula for pointing out the effective size difference.

m (GeV)	$\frac{Wu_k d_1}{\text{s.m.}}$	$\frac{Wt_1 b_k}{\text{s.m.}}$	$\frac{ZL_k L_1}{\text{s.m.}}$	$\frac{ZR_k R_1}{\text{s.m.}}$	m_t (GeV)	$\frac{Wt_k b_1}{\text{s.m.}}$	$\frac{Zt_k t_1 L}{\text{s.m.}}$	$\frac{Zt_k t_1 R}{\text{s.m.}}$
571.72	0.696	0.516	0.769	≈ 0	867.42	0.738	0.59	0.256
1514.91	0.551	0.409	0.617	≈ 0	1726.51	0.662	0.518	1.183
2491.81	0.512	0.387	0.576	≈ 0	2643.94	0.613	0.49	1.582
3472.06	0.495	0.378	0.559	≈ 0	3587.82	0.58	0.473	1.74
4450.21	0.485	0.374	0.549	≈ 0	4542.45	0.558	0.462	1.785
5423.85	0.479	0.371	0.543	≈ 0	5499.85	0.541	0.454	1.772
6391.3	0.475	0.369	0.538	≈ 0	6455.46	0.529	0.448	1.726
7351.12	0.472	0.368	0.535	≈ 0	7406.28	0.519	0.443	1.658
8302.01	0.469	0.367	0.533	≈ 0	8350.07	0.512	0.439	1.577

Table 2: The first nine recurrences for light quarks and the t, along with some of their couplings to light quarks and W and Z, compared to the standard model couplings.

2. Some of the couplings are quite large, and one certainly has to worry that loop-corrections will modify things.

3. To go beyond the 2 TeV level for the VEVs in this class of models requires even larger Yukawa and gauge couplings. But even if one does not worry about the size of the couplings, at some point, this program runs out of steam, because there is no way to get the appropriate gauge boson and t masses with real couplings.

4. Since there are lots of parameters, discrepancies at the percent level are interesting only if there are bounds that force them to go in a particular direction. That is probably the case for the couplings, WLL, ZLL and ZRR, which are systematically smaller than the tree-level standard model, presumably because some of the low-energy weak interactions arise from coupling of the light fermions to heavy gauge bosons. This model does not have "ideal delocalization" [18] in which these couplings vanish. My belief is that it is very difficult to implement this in any natural way so I find it heartening that at least these discrepancies can be made quite small by a suitable choice of parameters.

5. Searching even this small parameter space is greatly facilitated by having analytic expressions for many of the parameters, which is possible because of the translation invariance of the bulk. Very likely there are slightly warped models [27, 28, 29, 30, 31] nearby with even better correspondence to the standard model. And of course, there may be strongly warped solutions in very different regions of parameter space — I just do not know how to look for them efficiently.

6. The most interesting thing, I think, about this little exercise, is the unusual assignment of fermions to nodes in (4). While this makes perfect sense in the 4-d picture, it is not

obvious what it means for the interpretation of the deconstructed extra dimension. It is hard to see how color, for example, could be spread over such a construction.

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